

# Constraining Strong Baryon–Dark Matter Interactions with Primordial Nucleosynthesis and Cosmic Rays

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## Abstract

Self-interacting dark matter (SIDM) was introduced by Spergel & Steinhardt to address possible discrepancies between collisionless dark matter simulations and observations on scales of less than 1 Mpc. We examine the case in which dark matter particles not only have strong self-interactions but also have strong interactions with baryons. The presence of such interactions will have direct implications for nuclear and particle astrophysics. Among these are a change in the predicted abundances from big bang nucleosynthesis (BBN) and the flux of  $\gamma$ -rays produced by the decay of neutral pions which originate in collisions between dark matter and Galactic cosmic rays (CR). From these effects we constrain the strength of the baryon–dark matter interactions through the ratio of baryon - dark matter interaction cross section to dark matter mass,  $s$ . We find that BBN places a weak upper limit to this ratio  $\lesssim 10^8 \text{ cm}^2 \text{ g}^{-1}$ . CR-SIDM interactions, however, limit the possible DM-baryon cross section to  $\lesssim 5 \times 10^{-3} \text{ cm}^2 \text{ g}^{-1}$ ; this rules out an energy-independent interaction, but not one which falls with center-of-mass velocity as  $s \propto 1/v$  or steeper.

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# 1 Introduction

Compelling observational evidence suggest that matter in the universe is dominated by “cold dark matter” (CDM), the simplest model of which is one where dark matter particles have interaction strengths at or below the weak scale, and thus interact today only through gravity (*collisionless* dark matter). The collisionless cold dark matter model has been proven through analytic and numeric simulations to be very successful in explaining the large scale structure of the universe [1].

However, a series of observations of dark matter structures on the order of  $\lesssim 1$  Mpc (e.g., galaxy rotation curves [2, 3, 4], strong lensing from galaxy clusters [5], Tully-Fisher relation of spiral galaxies [6], bar stability in spirals [7, 8] and deficit of low-mass subhalos in the Local Group [9]) appear to contradict the prediction of collisionless CDM simulations that dark matter halos form high-density cores, favoring shallower, lower-density dark matter halo cores.

Spergel and Steinhardt [10] have proposed a variation of CDM in order to alleviate these discrepancies. In their proposed picture, dark matter particles interact with each other strongly and have a large scattering cross section, but negligible annihilation or dissipation.<sup>1</sup> Subsequent numerical simulations [11, 12] have shown that this collisional cold dark matter, or SIDM, does indeed predict halo cores in better agreement with the observations mentioned above, provided

$$0.5 < s < 6 \text{ cm}^2 \text{ g}^{-1}. \quad (1)$$

where the the parameter  $s = \sigma_{\text{SIDM}}/M_D$  is the dark matter–dark matter elastic scattering cross section over the dark matter particle mass. Potential difficulties for this model may arise at cluster scales [13, 5, 14, 15]; one way to address these would be a velocity-dependent cross section, the effect of which we will consider below.

If dark matter particles interact with each other through the strong force, then similar interactions would be expected between them and ordinary baryons as well, with cross sections of the same order [12]. Wandelt et al. [12] considered this case and found, perhaps surprisingly, that such interactions cannot be excluded by galaxy halo data nor from *direct* observation by space-borne cosmic-ray detectors, if dark matter particles have a mass larger than  $10^5$  GeV. Earlier papers [16] have also considered astrophysical constraints on SIDM both with and without interactions with baryons.

From the particle physics point of view, we choose a model-independent approach. The question we address is “If a massive neutral particle exists, how can we constrain its properties from

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<sup>1</sup> If the dark matter particle and antiparticle states were equally populated and annihilations were allowed, these would proceed vigorously, leaving a relic abundance today  $\Omega_D \sim 10^{-37} \text{ cm}^2 \langle \sigma_{\text{ann}} v/c \rangle^{-1}$  which would be negligible for the case of strong interactions. This implies that the SIDM particles either have strongly suppressed annihilation cross sections, or are asymmetrically populated (as is the case for, e.g., baryons).

astrophysical observations?”. We refrain from a detailed discussion of the naturalness of such a particle except to note how quickly model builders proposed several particle physics implementations (e.g. [17]) of self-interacting dark matter. Some promising candidates, the  $S0$  and low-mass strangelets, were then ruled out [12]. Nevertheless, there remain intriguing candidates for dark matter which lie in the regions of masses and baryon/self-interaction cross sections we consider, both within supersymmetric extensions of the standard model [18] or as possible relics from the QCD phase transition within the standard model [19]. Our results place constraints on the energy dependences of the cross sections in these models.

In this paper, we examine the effect of inelastic SIDM–baryon interactions. Specifically, we consider the impact of such interactions on big bang nucleosynthesis (BBN), and on the  $\gamma$ –rays produced by interactions of cosmic ray nuclei with dark matter particles. In BBN, the formation of the light elements deuterium,  $^3\text{He}$ ,  $^4\text{He}$ , and  $^7\text{Li}$ , depends on the time at which the weak interactions “freeze out,” the free neutron decay rate, and the temperature at which nucleosynthesis occurs, the latter being determined by the photo-dissociation of deuterium. If dark matter can interact strongly with baryons, it is possible that it can destroy light nuclei. Since deuterium has the smallest binding energy and determines the onset of nucleosynthesis, any dark matter deuterium dissociation will change where the onset of nucleosynthesis happens and thus provide a sensitive probe to baryon–dark matter interactions during BBN. We will see that the constraints provided by BBN are quite weak—in other words, for the parameter ranges of interest, SIDM has negligible effect on BBN despite the strength of the SIDM-baryon interaction.

The interaction of cosmic rays with interstellar matter leads to observable signatures in the form of  $\gamma$ –rays. Collisions between energetic cosmic ray protons and interstellar hydrogen can produce neutral pions, which in turn decay into two high-energy  $\gamma$ –rays; this emission is believed to dominate the observed flux of high-energy photons from the Galactic disk [20, 21, 22]. Including dark matter–baryon interactions allows for similar processes to occur between cosmic ray nuclei and ambient SIDM, thus possibly contributing to the  $\gamma$ –ray diffuse background. The  $\gamma$ –ray flux from CR-SIDM interactions will be a function of the assumed cross section for the interaction between dark and baryonic matter and can be calculated given the functional form of the Galactic cosmic ray flux and dark matter distribution. The result can then be compared with diffuse  $\gamma$ -ray observations from  $\gamma$ –ray telescopes (most recent of which are the EGRET observations, [21, 20]) and strong upper limits can be placed on the cross sections.

This paper is organized as follows. We will discuss the strong interacting dark matter cross section in §2. In §3 we determine the SIDM effect on BBN and its constraints. In §4 we compute the expected  $\gamma$ -ray flux from cosmic ray - SIDM interaction in the Milky Way and the resulting constraints on the cross section. Other possible constraints are noted in §5. Finally, our findings

are summarized and discussed in §6.

## 2 SIDM-Nucleon Interactions

We wish to characterize the inelastic interactions between baryonic nuclei and SIDM particles  $D$ . We will express the cross sections for inelastic interactions in terms of the elastic cross sections. The inelastic behavior is assumed to obey simple scalings with the masses of the reactants [10] and be energy-independent (but we will use the different constraints below to infer constraints on possible energy dependence of SIDM-baryon interactions).

We begin by considering the SIDM-nucleon interaction, with an energy-independent elastic scattering cross section  $\sigma_{DN}^{\text{elastic}}$ . This need not be the same as the SIDM self-interaction cross section, but was taken to be so in the work of [12]. One can then generalize this to determine the cross section for elastic SIDM-nucleus scattering. Following [16], we adopt the form

$$\sigma_{DA}^{\text{elastic}} = A^2 \left( \frac{\mu(A)}{\mu(p)} \right)^2 \sigma_{DN}^{\text{elastic}} \quad (2)$$

where  $A$  is the atomic mass number of the interacting nuclide, and  $\mu(A)$  is the reduced nucleus–dark matter mass. Finally, we will be interested in *inelastic* scattering, i.e., when internal degrees of freedom are excited in the baryonic component. We will write the (energy-independent) inelastic scattering cross section for channel  $i$  as

$$\sigma_{DA}^i = \alpha_i \sigma_{DA}^{\text{elastic}} \quad (3)$$

where  $\alpha_i$  encodes the relative strength of the inelastic interaction.

A typical strong interaction has a cross section,  $\sigma_{DN}^{\text{elastic}} \sim 1$  barn and it is assumed that a dark matter particle has a mass,  $M_D \gtrsim 1$  GeV. A useful parameter given these typical values for the cross section and mass is defined as

$$s = \sigma_{DN}^{\text{elastic}} / M_D \quad (4)$$

usually quoted in cgs units as  $\text{cm}^2 \text{ g}^{-1}$ . A useful conversion factor for particle physics units is  $1 \text{ cm}^2 \text{ g}^{-1} = 1.78 \text{ barn}/(\text{GeV}/c^2)$ .

## 3 Limiting SIDM with BBN

Primordial nucleosynthesis is the process by which the light elements, consisting mainly of deuterium,  $^3\text{He}$ ,  $^4\text{He}$ , and  $^7\text{Li}$ , were produced. Standard BBN is a one parameter theory, depending only on the baryon-to-photon ratio,  $\eta \equiv n_B/n_\gamma = 2.74 \times 10^{-8} \Omega_B h^2$ , where  $\Omega_B$  is the current baryon density relative to the critical density,  $\rho_{\text{crit}} = 3H_0^2/8\pi G$ , and  $H_0 = 100h \text{ km s}^{-1} \text{ Mpc}^{-1}$  is the present value of the Hubble parameter.

Before considering the effect of SIDM, it is useful to recall the basic physics of BBN. At the onset ( $T \gtrsim 1$  MeV;  $t \lesssim 1$  s) of the BBN epoch, the universe was dominated by relativistic particle species, initially consisting of photons, 3 species of light neutrinos, and  $e^\pm$  pairs; baryons were non-relativistic and essentially only  $n$  and  $p$  with no complex nuclei. At this stage, all of these particle species existed in thermal and chemical equilibrium, due to the interaction rates between particles being much faster than the expansion rate of the universe.

At a temperature around an MeV, and time about 1 s, the weak interactions, coupling the neutrinos to the other particle species and keeping the  $n$  and  $p$  in equilibrium, become ineffective (interaction rates less than  $H$ ). At this time the neutrinos fully decouple from the primordial plasma and the neutron abundance “freezes out”, relative to the proton abundance, with the exception of the occasional free neutron decay. Shortly after, the electrons and positrons start to become non-relativistic, and thus annihilate, producing photons, which in turn heats the plasma relative to the neutrinos. Trace amounts of the light nuclei, deuterium,  $^3\text{He}$ ,  $^4\text{He}$ , and  $^7\text{Li}$  exist, with abundances held below nuclear statistical equilibrium levels due to deuterium photo-dissociation by photons in the high-energy tail of the thermal distribution. This “deuterium bottleneck” persists until the universe cools to  $T \sim 0.07$  MeV, when energetic photons are sufficiently rare, at which point deuterium rapidly grows and then is burnt into the other light elements.

We now consider what role CDM plays in these series of events. Depending on the origin of the CDM and its coupling to Standard Model particles, CDM may or may not be in thermal equilibrium with the rest of the cosmic plasma. In the standard CDM scenario, the CDM only interacts weakly with itself and with Standard Model particles, and thus these interactions will freeze out well before the epoch of BBN. Then the dark matter cools rapidly due to Hubble expansion. Throughout BBN, the dark matter is non-relativistic and the Universe is radiation dominated. So in the standard case, dark matter plays no role at all.

What happens in the SIDM case? The dark matter is still non-relativistic, but now elastic collisions between baryons and dark matter will thermalize the dark matter at early times and bring  $T_D = T_\gamma$ . Dark matter–baryon decoupling will not occur until much later, and we note that this will potentially have interesting effects on structure formation<sup>2</sup>. However, for our study

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<sup>2</sup>The tight coupling of the baryon and photon fluids and the coupling between the dark matter and the baryons result in coupling the DM to the photons. This gives pressure to the dark matter fluid and changes the evolution of perturbations until the dark matter decouples from the baryons and then cools adiabatically. We conjecture that this effect leads to a suppression of fluctuation power on perturbation modes which entered the horizon before DM-baryon decoupling. Standard freeze-out arguments show that DM-baryon decoupling occurs at a redshift of  $z \sim 2 \times 10^4 \sigma_{DN}^{\text{elastic}} - (3/2)$ , where  $\sigma_{DN}^{\text{elastic}}$  is measured in barn. The characteristic mass scale of perturbations entering the horizon at this time is of order  $10^9 M_\odot (\sigma_{DN}^{\text{elastic}})^2$ . Note that this scaling now depends just on the cross section, not on the mass of the particle. This suggests that DM-baryon interactions reduce the dark matter perturbation power on galaxy scales — just those which motivated the study of interacting dark matter in the first place. We plan to study these effects in more detail in a future publication.

of BBN elastic scattering is not important except for thermalizing the dark matter. What does affect BBN are *inelastic* interactions with SIDM particles. Specifically, such collisions can impart internal energy to the nucleus and lead to its dissociation. Given that deuterium has the smallest binding energy and its photo-dissociation determines the onset of nucleosynthesis, we will focus on deuterium breakup; similar processes should occur for other nuclei but are not significant due to the higher binding energies.

We consider the reaction  $D+d \rightarrow D+n+p$  and neglect the rare, three-body, inverse reaction. In order for this interaction to happen the center-of-mass kinetic energy of the dark matter–deuteron system must be equivalent to at least the binding energy of the deuteron,  $B = 2.224$  MeV. We also assume that the break-up cross section is comparable to the scattering cross section,  $\sigma_{Dd}^{\text{bkp}} = \alpha_1 \sigma_{Dd}^{\text{elastic}}$ , where  $\alpha_1$  represents the efficiency of the break-up reaction and is of order unity. We relate the deuterium–dark matter scattering cross section to the proton–dark matter cross section as defined by eq. (2).

In making this interaction applicable for BBN, we need to compute the thermally averaged reaction rate,  $\langle \sigma v \rangle$ , where  $v$  is the relative velocity of the interacting particles. Since for our discussion, we consider a constant cross section with the appropriate energy threshold, this thermal average can be easily found.

$$\langle \sigma v \rangle = 5.221 \times 10^9 \alpha_1 \sigma_{DN}^{\text{elastic}} \frac{\left(1 + \frac{m_p}{M_D}\right)^2}{\left(1 + \frac{m_d}{M_D}\right)^{3/2}} \left(1 + \frac{B_9}{T_9}\right) T_9^{1/2} \exp\left(-\frac{B_9}{T_9}\right) \text{ cm}^3 \text{s}^{-1} \quad (5)$$

where  $m_p$  and  $m_d$  are the masses of the proton and deuteron respectively, and  $B_9$  and  $T_9$  are the binding energy of deuterium and the temperature expressed in units of  $10^9$  K.<sup>3</sup>

The addition of this dark matter reaction adds another channel for deuterium destruction in addition to photo-disintegration. Thus, it will take longer to build up enough deuterium so that nucleosynthesis can occur. To study the effects of the new interaction analytically, we will look at the quasi-static equilibrium (QSE) abundance of deuterium, as laid out by, e.g., [23, 24], and see how this changes when nucleosynthesis occurs. The abundance of deuterium will be most sensitive to the  $np \rightarrow d\gamma$ ,  $d\gamma \rightarrow np$  and  $dD \rightarrow Dnp$  reactions. The QSE abundance is then given by the ratio of sources to sinks.

$$Y_i^{\text{QSE}} \equiv n_i/n_B = \frac{\sum_{jkl} Y_j Y_k [jk \rightarrow il]}{\sum_{jkl} Y_l [il \rightarrow jk]} \Rightarrow Y_d = \frac{Y_n Y_p [np \rightarrow d\gamma]}{Y_\gamma [d\gamma \rightarrow np] + Y_D [dD \rightarrow Dnp]} \quad (6)$$

Here  $n_i$  is the  $i^{\text{th}}$  particle species number density,  $n_B$  is the baryon number density,  $Y_\gamma = 1/\eta$ , and

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<sup>3</sup> Throughout this analysis we assume that the dark matter remains non-relativistic during the epoch of BBN,  $M_D \gtrsim 10$  MeV. If dark matter is relativistic during this time, not only will the reaction rate take on a new form, but now the dark matter will contribute to the expansion of the universe.

$[jk \rightarrow il]$  are the reaction rates

$$[jk \rightarrow il] = N_A \rho_B \langle \sigma_{jk \rightarrow il} v \rangle \quad (7)$$

given by several compilations, where  $N_A = m_u^{-1}$  is Avogadro's constant, and  $\rho_B$  is the baryon mass density.

The forward and reverse reaction rates are related to each other by detailed balance. This tells us that the ratio of reverse to forward rates is given by the thermal equilibrium distributions of the interacting species. For  $np \rightleftharpoons \gamma d$  one finds the Saha expression

$$\frac{[d\gamma \rightarrow np]}{[np \rightarrow d\gamma]} = \frac{Y_n^{\text{EQ}} Y_p^{\text{EQ}}}{Y_d^{\text{EQ}} Y_\gamma^{\text{EQ}}} = 1.40 \times 10^5 T_9^{-3/2} \exp(-B_9/T_9) \quad (8)$$

Assuming dark matter has a mass  $M_D \gg m_p$ , the rate is given by the following expression.

$$Y_D[dD \rightarrow Dnp] = 5.221 \times 10^9 \alpha_1 \left( \frac{s}{\text{cm}^2 \text{ g}^{-1}} \right) \left( \frac{\Omega_D}{\Omega_B} \right) \left( \frac{\rho_B}{\text{g cm}^{-3}} \right) \left( 1 + \frac{B_9}{T_9} \right) T_9^{1/2} \exp\left(-\frac{B_9}{T_9}\right) \text{ s}^{-1} \quad (9)$$

The ratio of the two deuterium destruction rates is then

$$\frac{Y_D[dD \rightarrow Dnp]}{Y_\gamma[d\gamma \rightarrow np]} \approx 2.1 \times 10^{-9} \eta_{10} \alpha_1 \left( \frac{s}{\text{cm}^2 \text{ g}^{-1}} \right) \left( \frac{\Omega_D}{\Omega_B} \right) T_9 \quad (10)$$

assuming  $T_9 \ll B_9$  and where  $\eta_{10} = \eta/10^{-10}$ . For reasonable values of  $\alpha_1$  and  $s$ , we see that this ratio is tiny for the temperatures relevant to BBN. This simply reflects the smallness of  $\eta$ , i.e., the large photon-to-baryon (and photon-to-SIDM) ratio. The smallness of the SIDM contribution will lead to a very weak perturbation to the light elements.

Combining the results of eqs. (8), (9), and (10) and requiring  $\dot{Y}_d = 0$  gives the QSE deuterium abundance. The deuterium bottleneck ends, and the light elements are formed, when this increases to order unity, at the temperature

$$T_9 \approx \frac{25.81}{34.18 - \ln(\eta_{10}) - \frac{3}{2} \ln(T_9) + 2.1 \times 10^{-9} \eta_{10} \alpha_1 \left( \frac{s}{\text{cm}^2 \text{ g}^{-1}} \right) \left( \frac{\Omega_D}{\Omega_B} \right) T_9} \quad (11)$$

with  $T_9 \lesssim 1$ .

Thus one can see that with  $\frac{\Omega_D}{\Omega_B} \sim 10$ , one needs an extremely large cross section,  $\alpha_1 s \gtrsim 10^7 \text{ cm}^2 \text{ g}^{-1}$ , to cause a noticeable change in the light element predictions.  ${}^4\text{He}$  is the most sensitive to these changes. To derive a limit on SIDM, we will adopt a very generous observational lower bound  $Y_p \geq 0.230$  [25]. Using this, and a full numerical implementation of eq. (9) in the expanded BBN code discussed in [26], yields a limit  $\alpha_1 s \leq 1.263 \times 10^8 \text{ cm}^2 \text{ g}^{-1}$  for the case  $M_D \gg m_p$ . Given the rate dependence on dark matter mass and cosmological parameters we can derive an explicit form for the maximum cross section,

$$\alpha_1 s < 1.263 \times 10^9 \frac{\left(1 + \frac{m_d}{M_D}\right)^{3/2}}{\left(1 + \frac{m_p}{M_D}\right)^2} \left( \frac{\Omega_B}{\Omega_D} \right) \text{ cm}^2 \text{ g}^{-1} \quad (12)$$

for a baryon-to-photon ratio  $\eta_{10} = 5.8$  and requiring  $Y_p > 0.230$ . The constraints are only slightly different for different assumed  $\eta$  and helium abundances.

We see, therefore, that BBN cannot place strong constraints on dark matter mass or strong baryon cross section. In other words, energy-independent SIDM is completely compatible with light element constraints and BBN. The weakness of the constraint follows from the smallness of  $\eta$  and its imposition of the deuterium bottleneck which delays nuclear reactions until temperatures far below the binding energies of the light nuclei. This ensures that SIDM-deuteron interactions will be too weak to allow dissociation, and the lack of SIDM-nucleon bound states ensures that SIDM particles will not otherwise compete with nucleons as the light elements are built up.

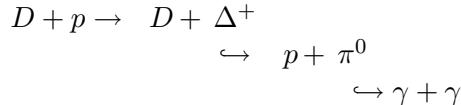
Having seen that SIDM-nucleon interactions are compatible with nucleosynthesis in the early universe, we now turn to a present-day consequence of these interactions, namely the production of  $\gamma$ -rays from SIDM collisions with cosmic rays.

## 4 Limiting SIDM with Cosmic Rays

The Galaxy is penetrated by a flux of nonthermal, relativistic nuclei, the cosmic rays. The bulk of these particles (i.e., those with  $E \lesssim 100$  TeV) are of Galactic origin, and have kinetic energies typically  $\sim 1$  GeV with a power law energy spectrum

$$\frac{dN}{dE} \propto E^{-(\gamma+1)} \quad (13)$$

where  $\gamma \sim 1.7$ . Inelastic collisions of cosmic ray nuclei (primarily protons) with SIDM particles are detectable in principle by the gamma-ray signature of the decay of the neutral pions produced in such interactions:



where  $D$  is the dark matter particle. This reaction is simply the SIDM analog of the usual pion production  $pp \rightarrow pp\pi^0 \rightarrow \gamma\gamma$  reaction, responsible for the bulk of the Galactic contribution to the  $\gamma$ -ray sky [27, 28, 21, 22].

The Galactic gamma-ray signature for this process is readily calculated. We assume axial symmetry for the cosmic-ray and SIDM spatial distributions, and a relevant dark matter–nucleon cross section  $\sigma_{DN}^{\text{inelastic}}$  independent of the incoming proton energy. The emissivity  $q_\gamma$  ( $\gamma$ -ray photons per volume per time) is then

$$q_\gamma(r, z) = 2 n_D(r, z) \sigma_{DN}^{\text{inelastic}} \Phi_p(r, z) . \quad (14)$$

Here, the factor of 2 is the number of photons per  $\pi^0$  decay,  $n_D(r, z)$  is the number density distribution of the dark matter in the Galaxy, and  $\Phi_p$  is the angle-integrated cosmic ray proton intensity

for all energies above the pion production threshold,

$$\Phi_p = 4\pi \int_{E_{\min}}^{\infty} \phi_{E,p} dE \quad (15)$$

with

$$E_{\min} = (m_p + m_{\pi^0}) \left[ 1 + \frac{m_{\pi^0}(2m_p + m_{\pi^0})}{2M_D(m_p + m_{\pi^0})} \right] \quad (16)$$

For  $M_D$  larger than  $\sim 1$  GeV, the second term in the brackets of eq. (16) is  $\ll 1$  and  $E_{\min}$  does not change appreciably with  $M_D$ , and we have  $E_{\min} \sim m_p + m_{\pi^0}$ . Using the parameter  $s = \sigma_{DN}^{\text{elastic}}/M_D$ , and the fact that with typical dark matter halo speeds  $v/c \sim 10^{-3} \ll 1$ , then  $\rho_D = M_D n_D$  and we can re-write eq. (14) as

$$q_{\gamma}(r, z) = 2\rho_D(r, z) \alpha_2 s \Phi_p(r, z). \quad (17)$$

The parameter  $\alpha_2$  as used in eq. (17), is the product of the ratio  $\sigma_{Dp \rightarrow DN\pi}/\sigma_{DN}^{\text{elastic}}$  and the branching fraction  $f_0$  for the decay of the  $\Delta^+$  to  $p + \pi^0$  rather than  $n + \pi^+$ . From isospin considerations it follows that  $f_0 = 2/3$ . In the absence of a detailed particle physics model for the interacting dark matter, it is natural to expect that the CR inelastic collision strength is related simply to the breakup collision strength by just this isospin factor,  $\alpha_2 = 2/3 \alpha_1$ , as both reactions are simply exciting internal degrees of freedom in the proton.

The resulting  $\gamma$ -ray intensity  $\phi_{\gamma, DN}$  in any given direction will then be the line integral of  $q_{\gamma}$  along the line of sight:

$$\phi_{\gamma, DN} = \frac{1}{4\pi} \int_{\text{l.o.s.}} q_{\gamma}(r, z) d\vec{\ell} \quad (18)$$

which, along a line of sight defined by a set of Galactic coordinates  $(\ell, b)$ , takes the form

$$\phi_{\gamma, DN} = \frac{1}{4\pi} \int_0^{\infty} q_{\gamma} \left( \sqrt{(R \cos b \cos \ell - L)^2 + (R \cos b \sin \ell)^2 + (R \sin b)^2}, R \sin b \right) dR \quad (19)$$

where  $L$  is the distance of the Sun from the Galactic center, and  $R$  is the heliocentric distance of a point of Galactocentric distance  $\vec{r}$ .

A first estimate of the expected  $\gamma$ -ray intensity from cosmic ray - dark matter interactions can be made if we approximate the spatial distribution of the CR flux by a step-function in space: Uniform and equal to the demodulated solar value inside an oblate ellipsoid with a semi-major axis of 10 kpc in the Galactic plane and 1 (5) kpc towards the Galactic poles, and zero outside this ellipsoid. In the limit  $M_D \rightarrow \infty$ , eqs. (15) and (16) give  $\Phi_{p,\infty} = 11.8 \text{ cm}^{-2} \text{ s}^{-1}$ , while if we relax our assumption of large  $M_D$ ,  $\Phi_p$  will scale as  $\Phi_p = \Phi_{p,\infty} \left[ 1 + \frac{m_{\pi^0}(2m_p + m_{\pi^0})}{2M_D(m_p + m_{\pi^0})} \right]^{-\gamma}$

In addition, we assume that the mass distribution of the dark matter halo can be described as a flattened, non-singular isothermal sphere [29],  $\rho(r, z) = \rho_c r_h^2 / (r_h^2 + r^2 + (z/q)^2)$ , where  $\rho_c$  is the central density,  $r_h$  the core radius and  $q$  the flattening of the halo. In the following calculations we have used  $q = 0.8$ ,  $\rho_c = 0.05 \text{ M}_{\odot} \text{ pc}^{-3}$  and a value of  $r_h = 5 \text{ kpc}$ , such that the total dark matter

Direction of Observation	SIDM expectation $\left(\frac{\alpha_2 s}{1\text{cm}^2/\text{g}}\right) \frac{\text{photons}}{\text{cm}^2 \text{sr s}}$	EGRET observation $\frac{\text{photons}}{\text{cm}^2 \text{sr s}}$	implied $\alpha_2 s$ upper limit
Galactic center	$2.1 \times 10^{-1}$	$5.5 \times 10^{-4}$	$2.6 \times 10^{-3}$
Galactic anti-center	$6.7 \times 10^{-3}$	$1.5 \times 10^{-4}$	$2.2 \times 10^{-2}$
Galactic poles	$2.5(12.) \times 10^{-3}$	$1.3 \times 10^{-5}$	$5.2(1.1) \times 10^{-3}$
M31	$9 \times 10^{-2}$	$< 8 \times 10^{-5}$	$9 \times 10^{-4}$

Table 1: Expected  $\gamma$ -ray intensities from SIDM-CR interactions and implied cross-section upper limits. The numbers in the parentheses indicate results for a semi-major axis of the cosmic ray halo ellipsoid in the direction perpendicular to the Galactic plane equal to 5 kpc, while other results refer to a semi-major axis of 1 kpc.

mass included in a radius of 100 kpc is  $\sim 10^{12} M_\odot$ . The distance of the Sun from the Galactic center was taken to be  $L \approx 8.5$  kpc.

We have calculated  $\phi_\gamma$  in this simple approximation for directions of observation towards the Galactic center ( $\ell = 0^\circ, b = 0^\circ$ ), towards the anti-center ( $\ell = 180^\circ, b = 0^\circ$ ) and perpendicular to the Galactic plane ( $b = 0^\circ$ ) and we have compared the results with the corresponding EGRET observational results from Hunter *et al* [21] and Sreekumar *et al* [20] to derive upper limits for  $\alpha_2 s$ . Our results for the limit of large  $M_D$  are shown in Table 1. The general result for the cross section upper limit will be given by

$$\alpha_2 s < \alpha_2 s_\infty \left[ 1 + \frac{m_{\pi^0}(2m_p + m_{\pi^0})}{2M_D(m_p + m_{\pi^0})} \right]^\gamma \quad (20)$$

where  $s_\infty$  is the value as  $M_D \rightarrow \infty$ .

Note that for a sufficiently flattened cosmic ray halo (1 kpc minor axis) the strongest constraint comes not from the direction where the  $\gamma$ -ray intensity is minimum (the Galactic poles), but from the observations towards the Galactic center, where the  $\gamma$ -ray intensity is maximum.

For the derivation of the upper limits presented above, the entire  $\gamma$ -ray flux detected by EGRET was attributed to pion decay from cosmic ray-SIDM interactions. This is a good assumption for the direction towards the Galactic poles where the contribution of the  $\gamma$ -rays from collisions of cosmic rays and interstellar baryons is not dominant. However, in the Galactic plane most of the observed intensity can be attributed to the presence of interstellar baryonic gas. Thus, the constraints in the directions of the Galactic center and anti-center could be made even more stringent than these conservative limits if one introduces a model for the emission from the interstellar medium.

We note that constraints of this type can be derived for external galaxies as well. In particular, we consider M31, which has the advantages of being nearby, and similar to the Milky Way. M31 has been searched for in *EGRET* sky maps, but no positive detection has been possible and an upper limit of  $1.6 \times 10^{-8} \text{ cm}^{-2} \text{ s}^{-1}$  has been placed instead on its diffuse  $\gamma$ -ray flux [30]. Given

the measured dark mass  $M_D^{\text{M31}} = 6 \times 10^{10} M_\odot$  in the inner 10 kpc of M31 [31], we predict a SIDM  $\gamma$ -ray flux of

$$\phi_\gamma^{\text{M31}} = 2 \frac{\Phi_p^{\text{M31}} \alpha_2 s M_D^{\text{M31}}}{4\pi R_{\text{M31}}^2} \quad (21)$$

where  $R_{\text{M31}} = 750$  kpc is the distance to M31. Following [32], we estimate the M31 cosmic ray flux  $\Phi_p^{\text{M31}}$  by assuming supernovae to be the site of cosmic ray acceleration, and thus taking the cosmic ray flux to be proportional to the supernova rate. This gives  $\Phi_p^{\text{M31}} = 0.45 \Phi_p^{\text{MW}}$ , which allows us to predict  $\Phi_\gamma^{\text{M31}} = 1.8 \times 10^{-5} \alpha_2 s \text{ cm}^{-2} \text{ s}^{-1}$  (where  $s$  is given in  $\text{cm}^2/\text{g}$ ). This must be lower than the observed limit, which implies a limit

$$\alpha_2 s^{\text{M31}} \lesssim 9 \times 10^{-4} \text{ cm}^2 \text{ g}^{-1} \quad (22)$$

Thus we see that the M31  $\gamma$ -ray limits are about an order of magnitude more constraining than those of the Milky Way; however, since they rely on more assumptions regarding the cosmic ray flux in M31, we regard the Milky Way limits as more secure.

Finally, it is very intriguing to note that a recent reanalysis [33] of the *EGRET* data finds evidence for a diffuse Galactic  $\gamma$ -ray halo. Specifically, a wavelet analysis was used to identify a  $\gamma$ -ray component at large angular scales, centered on the Galactic center, the intensity of which significantly exceeds the predictions of a model which includes known Galactic sources. One should bear in mind the error budget is large: the inferred halo intensity levels are uncertain to within a factor of  $\sim 2$ , and possible systematic effects remain. Also, it is entirely possible that the halo can be explained via inverse Compton scattering of background photons off of energetic cosmic ray electrons [22].

Nevertheless, it is tantalizing to interpret this diffuse halo in terms of SIDM-baryon interactions. The 100 MeV halo intensity levels toward the poles is quite uncertain but of order  $\sim 10^{-6} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ , i.e., about 10% of the full *EGRET* result. In terms of our model, this would amount to a *measurement* of  $\alpha_2 s$  at about 10% of our limit in Table 1, namely  $\alpha_2 s \sim 5(1) \times 10^{-4} \text{ cm}^2 \text{ g}^{-1}$ . Future  $\gamma$ -ray observatories such as *GLAST* can verify the existence of this diffuse Galactic halo, and will be in a position to measure its spectrum. If indeed the flux arises from  $\pi^0$  production in SIDM-baryon interactions, we would predict that the spectrum should reflect this origin, in particular, it should be a smooth continuum which shows the “pion bump” feature at  $E_\gamma = m_{\pi^0}/2 = 67.5$  MeV. In addition, similar halo emission should be visible around Local Group galaxies—not only M31, but also the Magellanic Clouds.

## 5 Other Possible Constraints

SIDM-nucleon interactions can have other effects as well. The possibility of DM-baryon interactions contributing to the gamma ray background provokes the question whether some fraction of the  $\sim 20$  unidentified EGRET sources at high Galactic latitudes could be due to dark matter concentrations in the Milky Way halo. A rough estimate shows that typical gamma-ray fluxes of these sources would be compatible with DM clumps of order  $10^3 M_\odot (R/1 \text{ kpc}) (1 \text{ cm}^2 \text{ g}^{-1} / \alpha_2 s)$ . While this is an intriguing speculation, the small number statistics of these sources makes distinguishing possible candidates from a background of approximately isotropically distributed AGN a difficult task. A more detailed study will therefore have to await the higher sensitivity and resolution data from future missions such as GLAST.

Another potential (but model-dependent) constraint comes from diffuse radiation. Consider the effect of SIDM particle passage through the interstellar medium of our Galaxy, most of which is hydrogen. Elastic collisions between SIDM particles and the protons occur, with a relative velocity given by  $v \sim 200 \text{ km/s} \sim 10^{-3}c$ , and thus a center-of-mass energy  $E_{\text{CM}} \simeq m_p v^2/2 = 0.2 \text{ keV}$ . Given that we expect  $m_p \ll M_D$ , it follows that a sizable fraction of this energy will be imparted to the recoiling proton. As a source of heating, this process does not violate observed constraints on the properties of molecular clouds [16]. However, if the scattering process leads to photon emission by the proton, this radiation might be observable. One can show that if each scattering event produces a photon with  $E_{\text{CM}}$ , then the resulting diffuse emission is at or above the level observed in soft X-rays. Of course, it need not be the case that this radiative scattering occurs, but given a particular SIDM model, this constraint should be investigated.

## 6 Discussion and Conclusions

We have considered constraints on strong interactions between dark matter particles and baryons. The big bang nucleosynthesis (BBN) constraints are weak, indicating that interacting dark matter, as it has been recently proposed, is fully consistent with BBN. However, very strong limits come from considering cosmic ray (CR) interactions with the dark matter. For dark matter masses  $M_D \gtrsim 1 \text{ GeV}$  the ratio  $s \equiv \sigma_{DN}/M_D$  is limited to be less than  $3.9 \times 10^{-3} \text{ cm}^2/\text{g}$ . Constraints from direct detection experiments such as the quantum calorimeter XQC [34] require  $M_D \gtrsim 10^5 \text{ GeV}$  [12, 35]. Our limit therefore constrains dark-matter baryon interactions to be less than a hundredth of the self-interaction strength which, simulations suggest, is required to affect the structure of galaxy halos within the self-interacting dark matter (SIDM) scenario. Any particle physics model of SIDM must explain this order of magnitude suppression of the dark matter-baryon interaction.

For the case in which the inelastic cross section is independent of energy, our results are sum-

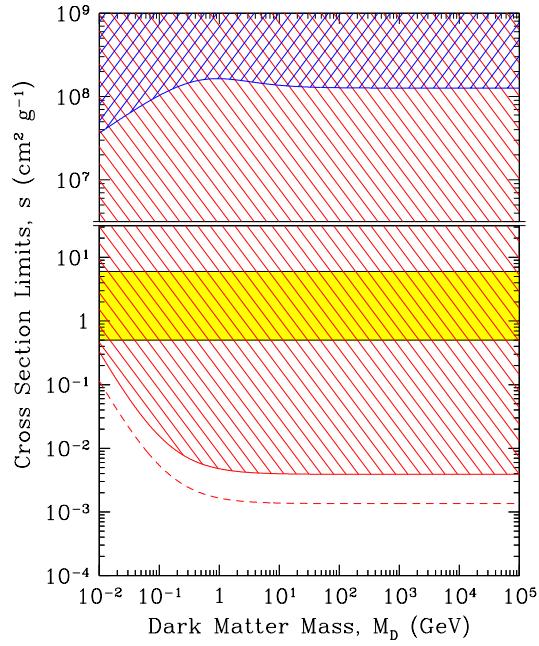


Figure 1: Limits on the cross section-to-mass ratio  $s$  as a function of SIDM mass. The values needed for galaxy halos (eq. 1) are shown in the shaded region. Excluded regions are hatched; single-hatched regions are excluded due to the inelastic interactions with Galactic cosmic rays; the cross-hatched regions are also excluded via BBN. The dashed curve shows the limit using M31  $\gamma$ -ray observations.

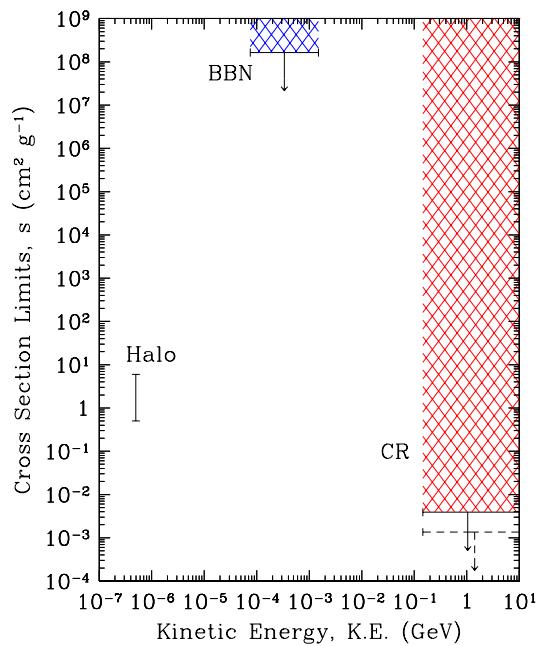


Figure 2: Limits on the cross section-to-mass ratio  $s$  as a function of center of mass energy, for the case in which  $M_D \gg m_p$ . The dashed curve shows the limit using M31  $\gamma$ -ray observations.

marized in Figure 1. It is of course possible, and indeed plausible, that SIDM-baryon interactions might be energy-dependent. Thus, we have displayed in Figure 2 the constraints on  $s$  as a function of the center-of-mass energy at which those constraints apply. Both figures show that energy-independent DM-baryon interactions are ruled out by the cosmic ray constraints, unless the ratio of DM-baryon scattering to DM self-scattering  $\alpha \lesssim 10^{-3}$ . On the other hand, if the interaction is energy-dependent, Figure 2 constrains the behavior. For example, if we assume that  $s \propto E_{CM}^n$ , where  $E_{CM}$  is the center of mass kinetic energy, then we find using our limits from cosmic rays coming from the Galactic center that  $n < -0.39$ . Using the M31 constraint steepens this to  $n < -0.47$ . In both cases, a  $1/v$  scaling (or stronger dropoff) is allowed. Thus, our  $\gamma$ -ray constraints provide important information about the energy dependence of a putative SIDM-baryon interaction. If such an energy dependent cross-section were characteristic of dark matter self-scattering as well, SIDM would be less interactive in objects with high velocity dispersion such as clusters of galaxies, possibly softening constraints on SIDM from the statistics of strongly lensed arcs [13].

We stress that our results do not rely on any detailed assumptions about the underlying particle physics model governing the interactions. In the context of a more detailed particle physics model for SIDM, one could also examine signatures of less generic processes, for example the generation of diffuse emission of soft X-rays from radiative scattering between SIDM and interstellar medium particles.

Finally, we note that elastic SIDM-baryon interactions thermalize the dark matter at early times. This may have potentially important consequences for structure formation. We plan to investigate this issue in more detail in a future publication.

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